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# The role of the pion cloud in the interpretation of the valence light-cone wavefunction of the nucleon

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## ABSTRACT

The pion cloud renormalises the light-cone wavefunction of the nucleon which is measured in hard, exclusive photon-nucleon reactions. We discuss the leading twist contributions to high-energy exclusive reactions taking into account both the pion cloud and perturbative QCD physics. The nucleon's electromagnetic form-factor at high  $Q^2$  is proportional to the bare nucleon probability  $Z$  and the cross-sections for hard (real at large angle or deeply virtual) Compton scattering are proportional to  $Z^2$ . Our present knowledge of the pion-nucleon system is consistent with  $Z = 0.7 \pm 0.2$ . If we apply just perturbative QCD to extract a light-cone wavefunction directly from these hard exclusive cross-sections, then the light-cone wavefunction that we extract measures the three valence quarks partially screened by the pion cloud of the nucleon. We discuss how this pion cloud renormalisation effect might be understood at the quark level in terms of the (in-)stability of the perturbative Dirac vacuum in low energy QCD.

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# 1 Introduction

The role of pions in light-cone QCD is a topic of much theoretical interest [1-7]. In this paper we explain how the pion cloud renormalises the valence light-cone wavefunction of the nucleon which is measured in the nucleon's electromagnetic form-factor at large  $Q^2$  and in hard Compton scattering at high energy. We also discuss how the pion cloud renormalisation of high-energy exclusive cross-sections might be understood in terms of dynamical symmetry breaking and the (in-)stability of the Dirac current-quark vacuum in low-energy QCD.

It is well known that the pion cloud, which is associated with chiral symmetry, plays an important role in the phenomenology of nucleon structure [8-11]. Pion cloud effects are present at all momentum scales. For example, in nuclear physics the process  $n \rightarrow p\pi^-$  offers a simple explanation of the long range part of the neutron's electric form-factor [12]. The pion cloud of the nucleon renormalises  $C = +1$  observables like the axial charge of the nucleon  $g_A^3$  as well as  $C = -1$  observables like the nucleon's anomalous magnetic moment  $\kappa_N$ . In high-energy deep inelastic scattering the process  $p \rightarrow n\pi^+$  generates a non-perturbative component in the nucleon's sea with an explicit anti-up, anti-down quark asymmetry [13]. This non-perturbative sea, together with the Pauli principle in the nucleon's wavefunction, explains in part [14] the violation of the Gottfried sum-rule [15] discovered by the NMC [16] at CERN – for a review see Ref. [17]. Mesonic effects also play some role in the explanation of the EMC nuclear effect [18, 19, 20].

In this paper we work on the light-cone and discuss the role of the pion cloud in the Fock expansion of the nucleon and in high-energy, exclusive photon-nucleon scattering. We concentrate on the nucleon's electromagnetic form-factor and on high-energy Compton scattering. We discuss how the light-cone wavefunctions which are measured in these processes should be interpreted in view of the new information we have learnt about the nucleon's internal structure in unpolarised and polarised (inclusive) deep inelastic scattering.

The structure of the paper is as follows. In Section 2 we start by discussing dynamical chiral symmetry breaking (D $\chi$ SB) in low energy QCD. The emphasis in this Section is on the key physical ideas and a supercritical phase transition as a possible explanation of the transition from current to constituent quarks. The picture of low energy QCD which emerges from such a phase transition is much like the Nambu-Jona-Lasinio model [11, 21]. In Section 3 we review the theory of how the pion cloud is included in the hadronic, light-cone Fock expansion of the nucleon and how the pion cloud of the nucleon contributes to deep inelastic scattering.

Our light-cone Fock expansion should provide a unified, self-consistent approach to both deep inelastic scattering and high-energy exclusive photon-nucleon reactions – the subject of Section 4. Whilst the pion cloud makes a leading twist contribution to deep inelastic structure functions, only the bare nucleon (leading hadronic Fock component) makes a leading twist contribution to the nucleon’s electromagnetic form-factor at large  $Q^2$  and to hard (real at large angle or deeply virtual) Compton scattering. At leading twist, the cross sections for high  $Q^2$ , elastic  $\gamma p \rightarrow p$  and hard Compton scattering are equal to the bare nucleon cross sections multiplied by the square of the bare nucleon probability  $Z$ , where  $Z = 0.7 \pm 0.2$  is determined from pion-nucleon physics. In our present theory of pion cloud effects in deep inelastic scattering [17], the parton model is defined with respect to the bare nucleon rather than directly with respect to the physical nucleon. The probability to find the physical nucleon in its three-quark leading Fock state is equal to the bare nucleon probability  $Z$  times the probability  $P_{3q}$  to find the leading Fock state in the bare nucleon. Section 4 concludes with a discussion how one might best extract information about the leading Fock state from present and future data on hard, exclusive processes. In Section 5 we explain how the pion cloud renormalisation of hard exclusive cross sections might be understood at the quark level in terms of the vacuum instability picture of  $D\chi$ SB in low-energy QCD outlined in Section 2. We discuss how one might construct a light-cone Fock expansion that includes both current [22, 23] and constituent [1, 7] quark degrees of freedom.

## 2 Dynamical chiral symmetry breaking in low energy QCD

QCD is asymptotically free. At large momentum transfer the running coupling  $\alpha_s(Q^2)$  decreases logarithmically with increasing  $Q^2$ . The expression for  $\alpha_s$  at one loop in perturbation theory is

$$\alpha_s(Q^2) = \frac{4\pi}{\beta_0 \ln \frac{Q^2}{\Lambda_{QCD}^2}}, \quad (1)$$

where  $\beta_0 = 11 - \frac{2}{3}N_f$  and  $N_f$  is the number of flavours. When  $Q^2$  is greater than about 2 GeV $^2$  the running coupling  $\alpha_s$  is small enough that one can apply perturbative QCD to calculate the short distance (or “hard”) part of a hadronic scattering process. The factorisation theorem [24, 25] then allows us to write hadronic cross sections as the convolution of “soft” parton distributions (in the case of inclusive deep inelastic scattering) or light-cone wavefunctions (in the

case of high energy exclusive reactions) with a “hard” scattering coefficient. The “soft” distributions contain all of the information about the structure of the target – the long-range bound state dynamics. They describe a flux of quark and gluon partons into the target independent “hard” part of the interaction.

Asymptotic freedom also gives us infrared slavery. The running coupling increases with decreasing resolution  $Q^2$ . Indeed, the perturbative expression for  $\alpha_s$  increases without bound if we let  $Q^2 \rightarrow \Lambda_{QCD}^2$  in Equ.(1). ( $\Lambda_{QCD}$  is the infrared Landau scale in QCD.) On the other hand, perturbation theory is derived assuming that the expansion parameter  $\frac{\alpha_s}{\pi} \ll 1$ . Physical arguments in non-perturbative QCD suggest that QCD may undergo a supercritical phase transition at a critical scale  $\lambda_c$  and that  $\alpha_s$  may “freeze” at the value  $\alpha_s^c = \alpha_s(\lambda_c)$  [26-31].

To understand what happens at this transition it is helpful to consider the analogous problem of a static, large- $Z$ , point nucleus in QED [30, 32, 33]. There the  $1s$  bound state level for the electron falls into the negative energy continuum at  $Z = 137$ . If we attempt to increase  $Z$  beyond 137 the point nucleus becomes a resonance: an electron moves from the Dirac vacuum to screen the supercritical charge which then decays to  $Z - 1$  with the emission of a positron.

First, let us consider QCD with just light quarks. If the quark itself were to acquire a supercritical charge at a critical scale  $\lambda_c$ , then it would not be able to decay into a positive energy bound state together with another quark with positive total energy because of energy momentum conservation. Instead, the Dirac vacuum itself would decay to a new supercritical vacuum state [30]. Since the vacuum is a scalar, this transition necessarily involves the formation of a scalar condensate which spontaneously breaks the (near perfect) chiral symmetry and yields the massive constituent quark quasi-particles of low energy QCD [26, 27, 31]. We call the phases at scales above and below the critical scale  $\lambda_c$  the Dirac and Landau phases of QCD respectively. Perturbative QCD is formulated entirely in the Dirac phase of QCD. The Dirac vacuum is a highly excited state at scales  $\mu \leq \lambda_c$  and one must re-quantise the fields with respect to the new ground state vacuum in the Landau phase of the theory. The Dirac quark of perturbative QCD would freeze out of the theory as a dynamical degree of freedom and the running coupling would freeze at  $\alpha_s(\lambda_c)$ , which is an infrared, unstable fixed point. The normal ordering mismatch between the zero point energies of the scalar vacua in the Dirac and Landau phases of QCD means that the quark in the low energy Landau phase feels a uniform, local potential which is manifest as the large mass of the constituent quark quasi-particle. The chiral dynamics of the Landau phase seem to be well described by the Nambu-Jona-Lasinio model

[11, 21, 26].

The freezing of  $\alpha_s$  is supported by the analysis of infrared induced, power-behaved contributions to hadronic event shapes in  $e^+e^-$  annihilation (for example, at LEP) [34] (see also [35]). Estimates of  $\alpha_s^c$  from these experiments range from 0.6 to 0.8. (It is also interesting to note that the pre- Sudakov effects [36] and pre- chiral symmetry analysis of the high  $Q^2$  behaviour of the nucleon’s Dirac form factor was consistent with freezing of  $\alpha_s$  at the (relatively small) value of 0.3 [37].)

Let us consider this scenario in more detail. In QCD with both light and heavy quarks there are two types of supercritical phase transition associated with a charged Dirac vacuum at large coupling  $\alpha_s$ : “static” transitions and “vacuum” transitions [30, 31, 38]. “Static” transitions involve the decay of a heavy quark  $q_h$  into a light quark  $q_l$  together with the formation of a positive energy ( $q_h\bar{q}_l$ ) meson bound state (like the decay of the large- $Z$  point nucleus in QED). These decays may, in part, be responsible for the confinement of heavy quarks. “Vacuum” transitions involve the decay of the fermionic vacuum for light-quarks from the Dirac into the Landau phase with the formation of a scalar condensate. Since the ( $q_h\bar{q}_l$ ) system has a higher reduced mass than the ( $q_l\bar{q}_l$ ) system, it follows that the static decay of a heavy quark would occur at a lower value of  $\alpha_s$ , or higher value of  $Q^2$ , than the vacuum transition involving light quarks.

To understand the confinement of light-quarks we make the simplifying hypothesis that the colour charge at  $\mu \leq \lambda_c$  is completely screened by the scalar condensate. Given that there is a supercritical vacuum transition, this hypothesis seems quite reasonable. If the Dirac quark acquires a supercritical colour charge, it then becomes a resonance in the negative energy continuum and decays without bound to yield the new vacuum state containing the scalar condensate and a dynamical colour charge that is “hidden” over any ultraviolet cutoff that we may choose to regularise the *Dirac* Fock space at  $\mu \leq \lambda_c$ . What remains is a massive spin  $\frac{1}{2}$  fermion which interacts with the condensate. (The formation of the scalar condensate and consequent mass generation stabilises the fermion vacuum [26].) This massive quasi-particle has a non-dynamical “passive” SU(3)-colour label which is manifest in the SU(3)-colour singlet hadronic wavefunctions of the constituent quark model. In this scenario the colour charge of an isolated quark is zero at distances greater than the critical radius  $r_c \sim \frac{1}{\lambda_c}$  and finite inside  $r \leq r_c$ . However, the colour charge is measured by a conserved vector current  $j^\mu$ , viz.  $D_\mu j^\mu = 0$  where  $D_\mu$  is the gauge covariant derivative. The total colour charge is conserved across the

critical radius  $r_c$ . If nature contained just one isolated quark, the critical scale would be infinite ( $\lambda_c \rightarrow \infty$ ) so that the colour charge would be completely screened. In this picture colour confinement is a local phenomenon. (This effect is essentially the same physics which prevents us from having massless charged particles in QED or in the Standard Model – see Refs.[38, 39].) For colour singlet hadrons (mesons and baryons) the nett colour charge is zero both inside and outside the critical radius  $r_c$  and the nucleon has finite size. The colour charge of a given quark is confined to scales  $\mu \geq \lambda_c$ . This scenario has phenomenological support in the many low energy properties of hadrons that can be described by the Nambu-Jona-Lasinio model, which includes chiral symmetry but not dynamical confinement [11, 27].

Clearly, this supercritical confinement scenario differs from the confinement that is found in pure gluodynamics and in quenched QCD on the lattice. The importance of chiral symmetry (light quarks) in hadron phenomenology poses an important question for lattice theorists: “Does this gluon induced confinement persist when we relax quenching and reduce the light quark mass to its physical value ?” The critical coupling for the Dirac vacuum to become unstable increases as we increase the light quark mass (so that fermion vacuum polarisation is suppressed). It seems reasonable that the confinement which is observed in pure gluodynamics may give way to fermion vacuum instability at some critical light quark mass. Whether this critical mass is above or below the physical light quark mass is an important question for future lattice calculations.

To conclude this section, we summarise how this physics offers a possible explanation of the transition from current to constituent quarks. In the Dirac phase of the theory (at a scale  $\mu > \lambda_c$  where  $\alpha_s < \alpha_s^c$ ) the fermionic degrees of freedom are the current quarks of perturbative QCD. In the low-energy Landau phase the fermionic degrees of freedom are massive, constituent quark quasi-particles. The valence quarks in a hadron are the minimal colour-singlet combination that enters the hadron’s wavefunction:  $q_i \bar{q}_i$  for a meson and  $\epsilon_{ijk} q_i q_j q_k$  for a baryon. (Here the subscript  $i$  refers to the colour of the quark.) The sea of quark-antiquark excitations in the Dirac vacuum condense in the vacuum transition from the Dirac phase to the Landau phase. As a result of this transition the valence current quarks acquire a large mass to become the valence constituent quarks of low energy QCD. The low energy quark-antiquark condensate is manifest in high-energy experiments. It gives us the infinite number of quark and antiquark partons which are observed in the unpolarised, deep inelastic structure function at small  $x$  and the need for a subtraction in the dispersion relation for the total  $\gamma p$  cross section [40].

Pions, as the lightest mass excitation of the condensate, should be included in the nucleon’s wavefunction. As we now discuss, they play an important role in high-energy photon-nucleon scattering.

### 3 The role of the pion cloud in deep inelastic scattering

Models of the nucleon which include chiral symmetry generally involve a bare nucleon and a pion cloud. The bare nucleon is defined by the SU(3) flavour, SU(2) spin wavefunction of the three valence constituent quark quasi-particles (in some confining potential).

Given this picture, the physical nucleon can be viewed on the light-cone (or in the infinite momentum frame) as the superposition of the bare nucleon and (in one-meson-approximation) two-particle meson-baryon Fock components [17, 41, 42, 43], viz.

$$|N(p)\rangle_{phys} = Z^{\frac{1}{2}} \left\{ |N(p)\rangle_{bare} + \sum_{M,B} \int dy dk_T^2 g_{0MBN} \phi_{MB}(y, k_T) |M, B(p, y, k_T)\rangle \right\}. \quad (2)$$

Here  $Z$  is the bare nucleon probability;  $\phi(y, k_T)$  is the probability amplitude to find the physical nucleon in a state  $|M, B(p, y, k_T)\rangle$  consisting of a meson  $M$  and a baryon  $B$  which carry light-cone momentum fractions  $yp_+$  and  $(1 - y)p_+$ , and transverse momentum  $k_T$  and  $-k_T$  respectively. Although we work in one-meson-approximation, we include higher order vertex corrections to the bare coupling  $g_{0MBN}$  and use the dressed hadronic coupling  $g_{MBN}^2 = Zg_{0MBN}^2$ . The probability to find the physical nucleon in a state consisting of the meson  $M$  and baryon  $B$  carrying  $y$  and  $(1 - y)$  percent of the physical nucleon’s light-cone momentum  $p_+$  is

$$f_{MB}(y) = g_{MBN}^2 \int dk_T^2 |\phi_{MB}(y, k_T)|^2. \quad (3)$$

Conservation of light-cone momentum  $p_+$ , expressed through the equation

$$f_{MB}(y) = f_{BM}(1 - y), \quad (4)$$

provides an important constraint on pion cloud models [42]. The number of mesons “in the cloud” is then

$$\langle n \rangle_{MB} = \langle n \rangle_{BM} = \int_0^1 dy f_{MB}(y) \quad (5)$$

and the bare nucleon probability is

$$Z = 1 - \langle n \rangle_{\pi N} - \langle n \rangle_{\pi \Delta}, \quad (6)$$

where we now restrict our attention to the pion cloud.

The Fock expansion in Equ.(2) can be used to study the effect of the pion cloud in deep inelastic scattering and high-energy exclusive reactions. At this point we note one important feature of light-cone perturbation theory: exchanged quanta are on-mass-shell and off-energy-shell. In deep inelastic scattering and deeply virtual Compton scattering the large momentum squared associated with the exchanged photon is  $Q^2 = q_T^2$  (strictly speaking  $s_{\gamma e'}$ ) and  $q_\mu q^\mu = 0$ .

Pion cloud contributions to deep inelastic scattering are calculated via the Sullivan process [44] shown in Fig. 1. In deep inelastic scattering we measure the inclusive cross section. The scattering of the hard photon from the two-particle meson-baryon Fock states is a leading twist effect. The struck hadron is “shattered” by the hard photon and, therefore, does not feel any final state  $\mathcal{O}(\frac{1}{Q^2})$  hadronic form-factor. When we include the Sullivan process, the parton distributions of the physical nucleon are obtained as the convolution of the pionic splitting functions  $f_{\pi N}(y)$  and  $f_{\pi\Delta}(y)$  with the parton distributions of the struck hadron, viz. [17]:

$$\begin{aligned} x(q \pm \bar{q})_{N,phys}(x, Q^2) = & Z x(q \pm \bar{q})_{N,bare}(x, Q^2) \\ & + x \int_x^1 \frac{dy}{y} f_{\pi N}(y) (q \pm \bar{q})_\pi \left( \frac{x}{y}, Q^2 \right) + x \int_x^1 \frac{dy}{y} f_{N\pi}(y) (q \pm \bar{q})_{N,bare} \left( \frac{x}{y}, Q^2 \right) \\ & + x \int_x^1 \frac{dy}{y} f_{\pi\Delta}(y) (q \pm \bar{q})_\pi \left( \frac{x}{y}, Q^2 \right) + x \int_x^1 \frac{dy}{y} f_{\Delta\pi}(y) (q \pm \bar{q})_{\Delta,bare} \left( \frac{x}{y}, Q^2 \right) \end{aligned} \quad (7)$$

and

$$\begin{aligned} xg_{N,phys}(x, Q^2) = & Z xg_{N,bare}(x, Q^2) \\ & + x \int_x^1 \frac{dy}{y} f_{\pi N}(y) g_\pi \left( \frac{x}{y}, Q^2 \right) + x \int_x^1 \frac{dy}{y} f_{N\pi}(y) g_{N,bare} \left( \frac{x}{y}, Q^2 \right) \\ & + x \int_x^1 \frac{dy}{y} f_{\pi\Delta}(y) g_\pi \left( \frac{x}{y}, Q^2 \right) + x \int_x^1 \frac{dy}{y} f_{\Delta\pi}(y) g_{\Delta,bare} \left( \frac{x}{y}, Q^2 \right). \end{aligned} \quad (8)$$

The pionic splitting functions are [17]:

$$f_{\pi N}(y) = \frac{3g_{\pi NN}^2}{16\pi^2} \int_0^\infty \frac{dk_T^2}{(1-y)y} \frac{\mathcal{F}_{\pi N}^2(s_{\pi N})}{(M_N^2 - s_{\pi N})^2} \left( \frac{k_T^2 + y^2 M_N^2}{1-y} \right) \quad (9)$$

$$\begin{aligned} f_{\pi\Delta}(y) = & \frac{4}{3} \frac{f_{\pi N\Delta}^2}{16m_\pi^2 \pi^2} \int_0^\infty \frac{dk_T^2}{(1-y)y} \frac{\mathcal{F}_{\pi\Delta}^2(s_{\pi\Delta})}{(M_N^2 - s_{\pi\Delta})^2} \\ & \frac{\left[ k_T^2 + (M_\Delta - (1-y)M_N) \right] \left[ k_T^2 + (M_\Delta + (1-y)M_N)^2 \right]^2}{6M_\Delta^2(1-y)^3}. \end{aligned} \quad (10)$$

Here  $\mathcal{F}_{\pi B}(s_{\pi B})$  is a  $\pi B N$  hadronic form-factor and  $s_{\pi B}$  is the invariant mass squared of the  $\pi B$  intermediate state

$$s_{\pi B}(k_T^2, y) = \frac{k_T^2 + m_\pi^2}{y} + \frac{k_T^2 + M_B^2}{1-y}. \quad (11)$$

In phenomenological analyses the form-factor  $\mathcal{F}_{\pi B}(s_{\pi B})$  is usually written using a dipole [41] or exponential [42] form. Since  $s_{\pi B}$  is invariant under  $(\pi(y) \leftrightarrow B(1-y))$  these pionic splitting functions satisfy the  $p_+$  conservation equation, Equ.(4). As a consistency check, note that the physical and bare nucleon distributions are both correctly normalised to the number of valence quarks,  $N_q$ , in the nucleon:

$$\int_0^1 dx (q - \bar{q})_{N,phys}(x, Q^2) = N_q \quad (12)$$

and

$$\int_0^1 dx (q - \bar{q})_{N,bare}(x, Q^2) = N_q \quad (13)$$

where  $(N_u = 2, N_d = 1)$  in the proton and  $(N_u = 1, N_d = 2)$  in the neutron. The process  $p \rightarrow n\pi^+$  generates an excess of anti-down quarks over anti-up quarks in the nucleon's wavefunction [13]. The Sullivan process, together with the Pauli principle in the nucleon's wavefunction, offers a simple explanation [14] of the violation of the Gottfried sum-rule observed by the NMC [16] – for a review see Ref. [17].

There is an important and subtle point to note from Equs.(7,8). The quark and gluon distributions of the physical nucleon which appear on the left hand side of Equs.(7,8) are the same quark and gluon distributions that appear in the operator product expansion analysis of deep inelastic scattering. The parton model distributions are defined via the factorisation theorem [25] as a flux of quark and gluon partons into the hard photon-parton scattering, which is described by the perturbative Wilson coefficients. *These parton model distributions are defined with respect to the bare baryon and pion and not directly with respect to the physical nucleon in this approach.* This result will be very important when we discuss high energy exclusive processes in Section 4.

Anticipating our discussion of exclusive reactions, we will need the value of the bare nucleon probability  $Z$ . This quantity is determined via Equs.(5,6) and (9,10) by the hardness of the pion-nucleon form-factor. In (equal time) nuclear physics applications the pion nucleon form-factor is commonly parametrised by a covariant monopole, viz.

$$\mathcal{F}_{\pi N}(k^2) = \frac{\Lambda_F^2 - m_\pi^2}{\Lambda_F^2 - k^2}. \quad (14)$$

Thomas and Holinde [45] (see also Holinde [46]) propose that  $\Lambda_F = 500 - 800\text{MeV}$  (with a preferred value of  $730\text{MeV}$ ) is consistent with pion nucleon phenomenology. This soft pion nucleon form-factor is now well accepted in the nuclear physics community [47-50]. A recent lattice calculation by Liu and collaborators [51] gives  $\Lambda_F = 750 \pm 140\text{MeV}$ , which is consistent

with the Thomas and Holinde result. Taking  $\Lambda_F = 650 \pm 150$  MeV in a covariant monopole is equivalent to a light-cone form-factor  $\mathcal{F}_{\pi N}(s_{\pi N})$  which corresponds to a bare nucleon probability

$$Z = 0.7 \pm 0.2 \quad (15)$$

where we assume equal hardness of  $\mathcal{F}_{\pi N}$  and  $\mathcal{F}_{\pi \Delta}$  and include an approximate 20% extra contribution from higher mass pseudoscalar and vector mesons [17]. In comparison, the value of  $Z$  which is calculated in the Cloudy Bag model is  $Z \simeq 0.5$  for a bag radius  $R = 0.8$  fm [52]. Some further renormalisation of the exclusive cross-section may come from Regge effects which become important when  $y \rightarrow 1$  in the branching process  $N \rightarrow N\pi$  and which are not included in the Fock expansion, Equ.(2) [53]. Given the theoretical uncertainties (working in one-meson-approximation, using renormalised couplings) it may be safer to consider the error on the bare nucleon probability as a uniform distribution instead of a normal distribution.

## 4 The role of the pion cloud in high-energy exclusive reactions

Following our discussion of pion cloud contributions to deep inelastic scattering, we now explain how the pion cloud is manifest in high-energy exclusive reactions. We focus on the nucleon's electromagnetic form-factor  $F_N(Q^2)$  at large  $Q^2 = q_T^2$  and high-energy Compton scattering.

When we analyse exclusive photon-nucleon reactions at high energy it is important to take into account both the pion cloud and also perturbative QCD physics. We first identify which diagrams contribute to these exclusive photon-nucleon reactions at leading twist when we work with the hadronic Fock expansion in Equ.(2). Following our discussion of deep inelastic scattering, we then apply the perturbative QCD factorisation of Brodsky and Lepage [23] to study the hard part of the exclusive reaction. In this way, we take into account both the physics of the Dirac and the Landau phases of QCD. (We shall discuss this point further in Section 5 below.)

The general rule when discussing pion cloud contributions to high-energy photon-nucleon exclusive processes is that diagrams which involve the flow of large momentum through a hadronic vertex are non-leading twist. This result follows because of the  $\mathcal{O}(\frac{1}{Q^2})$  denominator associated with the propagator of the struck hadron and also the  $\mathcal{F}_{\pi N}(\frac{1}{Q^2})$  hadronic form-factor suppression for the struck baryon (pion) to recombine with the spectator pion (baryon)

in flight to reconstruct the physical nucleon. The leading twist contribution to the nucleon's electromagnetic form-factor is

$$F_{N,phys}(Q^2) = Z F_{N,bare}(Q^2). \quad (16)$$

The bare nucleon probability  $Z$  renormalises the large  $Q^2$  part of the nucleon's electromagnetic form-factor. (For an explicit calculation of the higher twist pion cloud contributions to  $F_N(Q^2)$  within a particular pion cloud model, see Nikolaev et al.[54].) Note that  $F_{N,phys}(Q^2)$  is the form-factor that one would calculate in the Cloudy Bag model, including chiral symmetry, and  $F_{N,bare}(Q^2)$  is the form-factor that one would calculate retaining only the bare "MIT core" (without pions).

The hadronic Fock states in Equ.(2) which contribute to high-energy Compton scattering depend on the kinematics. Let  $q_i$  and  $q_f$  denote the momenta of the incident and emitted photons. Close to the forward direction in "near-real" ( $q_T \simeq 0$ ) Compton scattering, where  $(q_i - q_f)_\mu \ll \mathcal{O}(\Lambda_F)$ , there is an explicit pion cloud contribution to the physical cross section at leading twist. This contribution corresponds to diagrams where the incident and the emitted photons both couple to the same hadron in Compton scattering from the two-particle meson-baryon Fock component so that no large momentum flows through a hadronic vertex. In this "near-real", "near-forward" Compton scattering one has to consider explicit  $\gamma\pi \rightarrow \gamma\pi$ ,  $\gamma N \rightarrow \gamma N$  and  $\gamma\Delta \rightarrow \gamma\Delta$  contributions to the high-energy Compton scattering cross-section. In Compton scattering away from the forward direction the two-particle meson-baryon Fock states contribute only at higher twist because they involve the flow of large momentum through a hadronic vertex. Such hard processes are real Compton scattering at large angles and deeply virtual Compton scattering (dVCS), where the nucleon absorbs a large  $Q^2 = q_T^2 > \mathcal{O}(1\text{GeV}^2)$  photon and radiates a real photon into the final state. The cross-section for these hard Compton scattering processes is

$$d\sigma_{N,phys}(\text{hard CS}) = Z^2 d\sigma_{N,bare}(\text{hard CS}) \quad (17)$$

at leading twist. Having established which hadronic Fock components contribute to high-energy elastic photon-nucleon and Compton scattering, one can then apply the perturbative QCD analysis of Brodsky and Lepage [23] to  $F_N(Q^2)$  and to hard Compton scattering.

The counting rules tell us that the leading twist contribution to a hard, exclusive, photon-hadron scattering process is given by the hard photon scattering from the valence Fock component of the hadron involved in the hard scattering [24, 55, 56]. In their classic work on

high-energy exclusive reactions [23, 24], Brodsky and Lepage showed that perturbative QCD factorisation applies in these processes so that the leading twist part of the nucleon's electromagnetic form-factor can be written

$$F_{N,bare}(Q^2) = \int_0^1 [dx][dy] \Phi_{N,bare}^*(x_i, \mu) T_H(x_i, y_i, Q, \mu) \Phi_{N,bare}(y_i, \mu) \left[ 1 + \mathcal{O}\left(\frac{1}{Q}\right) \right] \quad (18)$$

where  $[dx] = dx_1 dx_2 dx_3 \delta(x_1 + x_2 + x_3 - 1)$  and  $\mu$  is the factorisation scale. Here  $\Phi_{N,bare}(x_i, \mu)$  is a valence light-cone wavefunction which describes the flux of valence quarks into the hard scattering, which is described by  $T_H(x_i, y_i, Q, \mu)$  and which may be calculated in perturbative QCD [23, 24, 36, 55, 57, 58]. This valence wavefunction is defined by the leading term in the Fock expansion of the nucleon on the light-cone in perturbative QCD — that is, in the Dirac phase of QCD.

Following Brodsky and Lepage [23], we introduce a partonic Fock expansion of the bare nucleon in perturbative QCD. If  $p_\mu$  and  $\lambda$  denote the nucleon's momentum and helicity respectively, then we write

$$|N(p, \lambda) \rangle_{bare} = \sum_{n, \lambda_i} \bar{\prod}_i \frac{dx_i}{\sqrt{x_i}} \frac{d^2 k_{T,i}}{16\pi^3} \psi_{n/N}(x_i, k_{T,i}, \lambda_i) |n : x_i p_+, \vec{p}_T + \vec{k}_{T,i}, \lambda_i \rangle \quad (19)$$

where

$$\bar{\prod}_i dx_i d^2 k_{T,i} = \prod_i dx_i \delta(\sum_i x_i - 1) d^2 k_{T,i} 16\pi^3 \delta(\sum_i k_{T,i}) \delta(\sum_i \lambda_i - \lambda). \quad (20)$$

In Equ.(19),  $\psi_{n/N}(x_i, k_{T,i}, \lambda_i)$  is the amplitude for finding the (bare) nucleon in the specific Fock state  $n$  consisting of partons with momenta  $(x_i p_+, x_i p_T + k_{T,i})$  and helicities  $\lambda_i$ . Integrating over the parton's transverse momentum  $k_{T,i} < \mu$ , the light-cone wavefunction

$$\Phi(x_i, \mu) = \int^\mu \frac{d^2 k_T}{16\pi^3} \psi_{n/N}(x_i, k_{T,i}, \lambda_i) \quad (21)$$

describes a flux of partons, collinear up to  $k_T < \mu$ , into the hard scattering described by  $T_H(x_i, y_i, Q, \mu)$ .

In the parton model these light-cone wavefunctions are defined with respect to the bare nucleon. To see this, consider the (higher-twist) two-particle meson-baryon Fock state contribution to  $F_N(Q^2)$ . The hard scattering takes place either on the baryon with the pion in flight or on the pion with the baryon in flight. The particle in flight does not participate in the hard scattering process, whence the factorisation theorem [23, 24] tells us to use the light-cone wavefunction of the bare nucleon in the perturbative QCD part of the analysis of high-energy exclusive scattering. The probability to find the physical nucleon in its leading three-quark,

valence Fock state is equal to the bare nucleon probability  $Z$  [9, 52] times the probability

$$P_{3q}(Q) = \int_0^1 [dx] \Phi_{bare}^*(x_i, Q) \Phi_{bare}(x_i, Q) \quad (22)$$

to find the leading Fock state in the bare nucleon [23, 24]. The shape of  $\Phi_{bare}(x_i, Q)$  determines how the nucleon's light-cone momentum is distributed among the three valence quarks in the leading Fock component. Given the probability interpretation of these parton model light-cone wavefunctions, they can also be used to calculate the structure function of the bare nucleon in deep inelastic scattering [24]

$$F_{2,bare}(x, Q^2) = x \sum_{a=q,g} \int_x^1 \frac{dy}{y} \sum_{n,\lambda_i} \int \prod_i \frac{dz_i d\vec{k}_{T,i}}{16\pi^3} |\psi_{n/N}^{(Q)}(z_i, \vec{k}_{T,i}, \lambda_i)|^2 \sum_{b=a} \delta(z_b - y) C^a\left(\frac{x}{y}, \alpha_s\right), \quad (23)$$

where we have included the quark charges into the perturbative coefficients  $C^a\left(\frac{x}{y}, \alpha_s\right)$ .

The light-cone wavefunction of the leading Fock component in the nucleon is calculated from the Fourier transform of the vacuum to nucleon matrix element of the nucleon interpolating operator  $\epsilon^{ijk} u_\alpha^i(z_1) u_\beta^j(z_2) d_\gamma^k(z_3)$ . (We refer to Chernyak and Zhitnitsky [59] for technical details of this calculation). The structure of the nucleon enters the calculation of  $\Phi(x_i, Q)$  in the light-cone matrix element

$$\langle vac | u_\alpha^i(z_1) u_\beta^j(z_2) d_\gamma^k(z_3) | p \rangle \epsilon^{ijk}, \quad (24)$$

which plays an analogous role to the light-cone correlation function in deep inelastic scattering. Modulo flavour and spin labels, the light-cone wavefunction of the leading Fock component can be expanded in terms of the set of orthogonal Appell polynomials  $A_n(x_i)$  [23, 59], viz.

$$\Phi(x_i, Q) = f_N(Q^2) \phi_{as}(x_i) \sum_n f_n(Q^2) a_n A_n(x_i). \quad (25)$$

Here  $\phi_{as} = 120x_1x_2x_3$  is the asymptotic, free-quark-model wavefunction,  $a_n$  are the expansion coefficients, and  $f_N$  and  $f_n$  carry the anomalous dimension of the nucleon interpolating operator – that is, they describe the  $Q^2$  dependence of the light-cone wavefunction.

The normalisation of any theoretical prediction of the light-cone wavefunction  $\Phi(x_i, Q)$  depends on the input that one uses for the proton state  $|N(p, \lambda) \rangle$  and the quark operators in Equ.(24). (There is no a-priori normalisation of a Bethe-Salpeter amplitude.) For example, consider the Cloudy Bag model. The model prediction of the light-cone wavefunction of the leading Fock component in the physical nucleon is equal to  $\sqrt{Z}$  times the three-quark, valence wavefunction of the bare (“MIT core”) nucleon, which is the light-cone wavefunction used in

the parton model. (The wavefunction of the physical nucleon measures the three valence quarks partially screened by the pion cloud.) For a bag radius  $R = 0.8\text{fm}$  [52] this means that the pion cloud renormalises the MIT bag model exclusive cross-section by a factor of four ! Given a bare nucleon probability  $Z = 0.7 \pm 0.2$  it is clearly very important to quantify the extent to which pion corrections (chiral symmetry) are included in any given model calculation of  $\Phi(x_i, Q)$  before comparing with data.

Quenched lattice calculations of  $\Phi(x_i, Q)$  [61] include some but not all pion loop effects [62, 63]. The situation is somewhat better in QCD sum-rule calculations [59, 64]. However, it is important to keep in mind that whereas the pion cloud renormalises the nucleon mass by about 30% [9], it can renormalise the cross section for hard Compton scattering by up to a factor of four. QCD sum-rule predictions for the unpolarised, hard Compton scattering cross-section typically differ by a factor of 2-3 [57] and, therefore, should not be distinguished by comparison with the absolute, measured cross-section alone.

The shape of the spin-independent, parton-model wavefunction for the three valence quarks without pionic dressing ( $\phi_{as}(x_i) \sum_n f_n(Q^2) a_n A_n(x_i)$  in Equ.(25)) is best determined from experiment by a maximum likelihood fit to  $Z$  independent ratios of unpolarised, hard, exclusive observables such as:

$$R_1 = \frac{d\sigma(\text{dVCS}, \theta)}{d\sigma(\text{dVCS}, \theta = 90^\circ)} \quad (26)$$

$$R_2 = \frac{d\sigma(\text{dVCS}, \theta)}{F_N^2(Q^2)}. \quad (27)$$

In performing these fits, it is important to make sure that the observables in the numerator and denominator are described self-consistently using the same factorisation scheme and also at the same factorisation scale (to eliminate factors of  $f_N$ ). The sensitivity of exclusive cross-sections to the normalisation of light-cone wavefunctions has been stressed previously by Hyer [58]. Given the fairly large uncertainty in the bare nucleon probability, one should compare the predictions of lattice [61], QCD sum-rule [59, 64] and diquark [60] models for the shape of light-cone wavefunctions with  $Z$  independent ratios of hard, exclusive cross-sections instead of the absolute cross-sections.

There has been much theoretical work in recent years aimed at understanding the EMC spin effect [65] in polarised deep inelastic scattering. Two important theoretical results are the role of the axial anomaly in spin dependent processes [66], and possible contributions from QCD background fields [67, 68]. The anomaly can induce a contact interaction between a

hard photon and the background field (on non-perturbative vacuum) [67, 69], which has no Fock representation in perturbation theory. This interaction is, in general, leading twist and has the potential to screen the spin of the quarks at large  $x$  in the spin dependent structure function  $g_1$  [67]. The phenomenology of this effect is that any  $C = +1$  spin observable is, in principle, subject to a significant violation of Zweig’s rule. This Zweig’s rule violation has the potential to modify the shape as well as the normalisation of the spin dependent light-cone wavfunction which is measured in hard, exclusive scattering. How might one try to isolate such a Zweig’s rule violation ? Experimentally, this requires a comparison between the light-cone wavefunction which is measured in  $C = +1$  and  $C = -1$  spin observables. (Such  $C = -1$  spin observables are found in parity-violating, hard Compton scattering.) To isolate the possible Zweig’s rule violation in theoretical calculations, one needs to introduce a switch in the QCD sum-rule calculation with which to turn on and off the effect of the anomaly — possibly along the lines suggested by Narison, Shore and Veneziano in polarised deep inelastic scattering [70].

## 5 $D\chi$ SB, confinement and light-cone QCD

In Sections 3 and 4 we introduced a two-stage Fock expansion to describe high-energy exclusive reactions in light-cone QCD. At the first stage we introduced pions explicitly and at the second stage we introduced (perturbative) quark and gluon degrees of freedom. It is worthwhile to stop and ask whether one could expand the physical nucleon directly in terms of perturbative quark and gluon Fock components and not lose any of the physics.

In light-cone perturbation theory all particles are on-mass-shell and potentially off-energy-shell. For a given particle

$$k_+ = \frac{m^2 + k_T^2}{k_-} \geq 0, \quad (28)$$

where the equality holds only for massless particles. When we quantise perturbative QCD on the light-cone the current quark has a well defined mass. We can set  $k^2 = m^2$  so that  $k_+ > 0$  and successfully apply light-cone perturbation theory [22]. Confinement becomes important at large coupling — at which point the Dirac quark has no well defined mass shell in (equal-time) QCD. If the Dirac quark becomes a supercritical resonance, then  $k_+$  develops an imaginary part and the  $k_+ \geq 0$  constraint no longer applies. (Even the real part of  $k_+$  can go negative.) Conventional light-cone perturbation theory [22] does not apply in an unstable vacuum. Motivated by our discussion in Section 2, this problem is cured if we work with massive constituent quarks

with non-dynamical colour, such as in the Nambu-Jona-Lasinio model [11, 21]. Of course, we need to remember that colour is confined. This means that the bare nucleon (composed of constituent quarks) is the intermediate step between the physical nucleon and perturbative QCD in the Fock expansion of the nucleon. We do not introduce a perturbative QCD Fock expansion of the constituent quark as the intermediate step because the constituent quark has no meaning when considered in isolation. (There are no free quarks.)

Finally, we outline how the chiral symmetry renormalisation of exclusive cross-sections might be understood at the quark level in terms of an unstable Dirac vacuum. Fradkin and collaborators [71] have derived the transition probabilities for exclusive reactions in QED with an unstable vacuum, which they include via an external field. Here, one needs to consider contributions to exclusive cross-sections where the vacuum “decays” during the exclusive scattering with the emission of  $e^-e^+$  pairs. (Remember that the  $k_+ \geq 0$  constraint no longer applies once we turn on vacuum instability.) It follows that

$$p_V = | \langle vac(out) | vac(in) \rangle |^2 < 1 \quad (29)$$

in an unstable vacuum. The creation (annihilation) operator for out-state electrons in QED with an unstable vacuum is a non-trivial, linear superposition of the creation (annihilation) operator for in-state electrons and the annihilation (creation) operator for in-state positrons [71]. Vacuum instability mixes the Fock components of the in- and out- state wavefunctions.

Turning now to QCD, this effect is clearly not important at short distances where  $\alpha_s$  is small – the “hard” part of the exclusive reaction. It is potentially very important for our interpretation of the “soft” light-cone wavefunctions if the Dirac quark is allowed to undergo a supercritical decay between light-cone times  $\tau \rightarrow -\infty$  and  $\tau \rightarrow +\infty$  in the exclusive reaction. In this case, what we call a quark in the in-state is modified by the charged vacuum en route to what we call a quark in the out-state of the scattering process. The exclusive cross-section is then suppressed with respect to the cross section that we would predict if we assume that the Dirac vacuum is stable at all scales. This suppression is identified with the bare nucleon probability in the vacuum instability mechanism for dynamical symmetry breaking outlined in Section 2 (which suggests a possible approach how one might calculate  $Z$  in an eventual solution to non-perturbative QCD). The supercritical decay is “static” for a heavy quark and “vacuum” for a light quark. The flavour-SU(N) and spin-SU(2) valence wavefunction of a given hadron determines the relative importance of “static” and “vacuum” transitions at large coupling  $\alpha_s$  and, therefore, the hadron dependence of the bare hadron probability  $Z$ .

## 6 Conclusions

Recent experiments in inclusive, deep inelastic scattering have shed new information on the role of the pion cloud (and possible background field effects) in the structure of the nucleon. These discoveries are also important to our interpretation of the light-cone wavefunctions which are measured in hard, exclusive scattering. The pion cloud renormalises the cross-sections for high-energy exclusive photon-nucleon processes. The nucleon's electromagnetic form-factor  $F_N(Q^2)$  at large  $Q^2$  is renormalised by the bare nucleon probability  $Z$  and the cross section for hard Compton scattering is renormalised by  $Z^2$ . The nucleon's valence light-cone wavefunction in the parton model is defined with respect to the bare nucleon (the MIT core in the Cloudy Bag model of the nucleon). The probability to find the physical nucleon in its leading three-quark, valence Fock state is equal to the bare nucleon probability  $Z$  [9, 52] times the probability  $P_{3q}$  to find the leading Fock state in the bare nucleon [23, 24]. Given the fairly large uncertainty on the value of  $Z$  (nearly a factor of two), it is important to compare the predictions of various models (lattice, QCD sum-rules, diquark) with  $Z$  independent ratios of hard, exclusive observables rather than the absolute cross-sections. If we apply just perturbative QCD to extract a light-cone wavefunction directly from the cross-sections for hard real and deeply virtual Compton scattering data, then the light-cone wavefunction that we extract has the interpretation that it measures the three valence quarks partially screened by the pion cloud of the nucleon. Dirac vacuum instability in low-energy QCD offers a possible quark level explanation of the pion cloud renormalisation of hard, exclusive photon-nucleon cross-sections.

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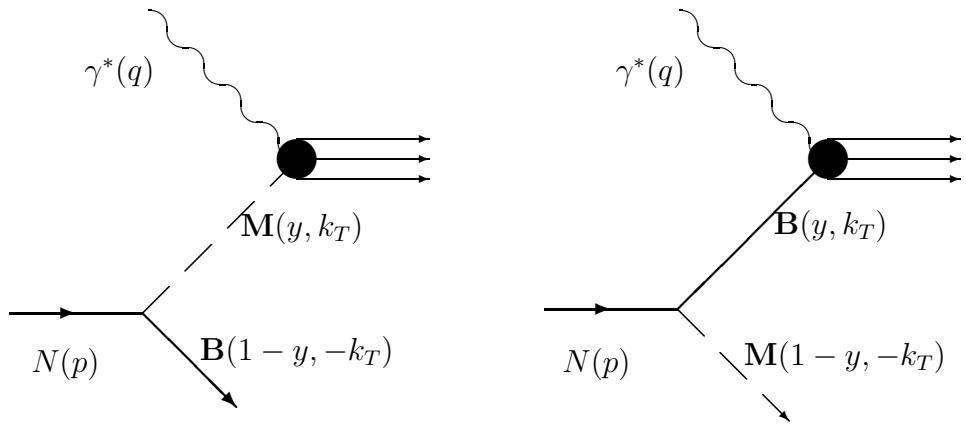


Fig.1: Two-particle meson-baryon contributions to deep inelastic scattering